

# Charmonium suppression at RHIC: Signature of a strongly-interacting QGP, not a weakly interacting

Binoy Krishna Patra<sup>1\*</sup> and Vineet Agotiya<sup>1</sup>

<sup>1</sup> *Department of Physics, Indian Institute of Technology Roorkee, Roorkee-247 667, India*

Following a recent work on equation of state for strongly interacting quark-gluon plasma [1], we revisited the equation of state by incorporating the nonperturbative effects in the deconfined plasma phase. Our results on thermodynamic observables *viz.* pressure, energy density, speed of sound etc. nicely fit with the lattice equation of state for gluon, massless and as well *massive* flavored plasma. Motivated by this agreement with lattice results, we have employed our equation of state to estimate the quarkonium suppression in an expanding, dissipative strongly interacting QGP produced in relativistic heavy-ion collisions and our prediction matches exactly with the recent PHENIX data on the centrality dependence of  $J/\psi$  suppression in Au+Au collisions at BNL RHIC. We have also predicted for the  $\Upsilon$  suppression in Pb+Pb collisions at LHC energy which could be tested in the ALICE experiments at CERN LHC.

**KEYWORDS:** Equation of State, Strongly Coupled Plasma, Heavy Quark Potential, String Tension, Dissociation Temperature

**PACS numbers:** 25.75.-q; 24.85.+p; 12.38.Mh ; 12.38.Gc, 05.70.Ce, 25.75.+r, 52.25.Kn

## I. INTRODUCTION

Quantum chromodynamics (QCD) at high temperature is believed to be in quark gluon plasma (QGP) phase, whereby color charges are screened rather than confined. Asymptotic freedom of non-abelian gauge theories insures that for high enough temperature  $T \gg \Lambda_{QCD}$ , QGP is weakly coupled with dressed quarks and gluons behaving as quasi-particles near the ideal gas limit.

In relativistic heavy ion collisions at RHIC, two novel phenomena - cone and ridge, not present in  $pp$  or  $d+Au$  collisions, were observed [2]. Quark-gluon plasma which is produced in heavy ion collisions is not an ideal gas of quarks and gluons, it is rather a liquid having very low shear viscosity to entropy density ( $\eta/s$ ) ratio [3–6]. This strongly suggest that QGP may lie in the non-perturbative domain of QCD which is very hard to address both analytically and computationally. Similar conclusion about QGP have been reached from recent lattice studies which predict that the equation of state (EoS) is interacting even at  $T \sim 4T_c$  [7–10]. Why QGP is strongly interacting in this temperature range and what is its meaning are not very well understood till the date. Since then, attempts have been made to understand strongly interacting nature of QGP and its small  $\eta/s$  ratio, argued from large elliptic flow observed in RHIC data. Similar conclusions about the near perfect fluidity of QGP have been reached

from the AdS/CFT studies [6], spectral functions and transport coefficients in lattice QCD [11] and studies based on classical strongly coupled plasmas [12, 13].

There are several attempts to model the EoS of such strongly-interacting matter *viz.* bag model, confinement models, quasi-particle models, strongly interacting quark gluon plasma[1, 14] etc. In the bag model [15], QGP is treated as a big hadron with large number of partons interacting weakly but confined by the bag wall. Further inclusions of glue balls or hadrons improve the predictions. The confinement models are the extension of bag model with smooth potential like Cornell potential [16] etc. with a better predictions. There is an interesting attempt [17, 18] to determine the EoS of interacting quarks and gluons upto  $O(g^5)$  which has been further improved upon [19, 20] by incorporating the contributions from the nonperturbative scales *viz.*  $gT$  and  $g^2T$  up to  $O(g^6 \ln(1/g))$ [21]. On the other hand, a semiclassical approach aimed to study the bulk properties of QGP automatically incorporate hard thermal loop (HTL) effects[22, 23] where the non-perturbative features manifested as effective mean color fields [24].

There are different versions of quasi-particle models [25] where equation of state was derived with temperature dependent parton masses and bag constant[26, 27], with effective degrees of freedom [28], etc. All of them claim to explain lattice results, either by adjusting free parameters in the model or by taking lattice data on one of the thermodynamic quantity as an input and predicting other quantities. However, physical picture of quasi-particle model and the origin of various temperature dependent quantities are not clear yet [29]. In strongly interacting QGP [30–32], one considers all possible hadrons even at  $T > T_c$  and try to explain non-ideal behavior of QGP near  $T_c$ . Recently, an equation of state for strongly-coupled plasma has been inferred by utilizing the understanding from strongly coupled QED plasma [33] which fits lattice data well. To be honest, deep and comprehensive

---

\*Electronic address: binoyfph@iitr.ernet.in

understanding is still missing and future investigations (both theoretical and forthcoming experiments at LHC) may throw more light on this very exciting and complex issue.

A suppressed yield of quarkonium in the dilepton spectrum, measured in experiments [34, 35] was proposed as a signature of QGP formation. Understanding the experimental data, however, turned out to be more complicated because the suppression pattern seen is not only due to the hot medium effects of screening [11, 36], but more likely due to interplay of different interactions *viz.* cold nuclear matter [37] effect, gluonic dissociation [38] as well as those of recombination [39] of  $c\bar{c}$  pairs etc. In order to disentangle these different effects we must know the properties of quarkonium in medium and their dissociation. There are two main lines of theoretical studies to determine quarkonium spectral functions at finite temperature: potential models [40]-[47] which have been widely used to study quarkonium, but their applicability at finite temperature is still under scrutiny and lattice QCD [11, 48] which provides the most straightforward way to determine spectral functions, but the results suffer from discretization effects and statistical errors, and thus are still inconclusive.

The short and intermediate distance properties of the heavy quark interaction are important for the understanding of in-medium modifications of the heavy quark bound states whereas the large distance behavior plays a crucial role in understanding the bulk properties of the QCD plasma phase [45]. In the study of bulk thermodynamic quantities *e.g.* pressure, energy density etc., deviations from perturbative calculations and ideal gas behavior are found at temperatures much larger than the deconfinement temperature which is further supported by robust collective flows, strong jet and charm quenching, and charm flow, in RHIC data. This calls for a quantitative nonperturbative calculations. Since the phase transition in full QCD appears as a crossover rather than a true phase transition [49], it is reasonable to assume that the string tension does not vanish abruptly above  $T_c$ . So one should study its effects on the behavior of quarkonia in a hot QCD medium. Recently we [50, 51] had considered this potentially interesting issue by correcting the full Cornell potential with a dielectric function embodying the effects of the deconfined medium and not only its Coulomb part as usually done in the literature. This led to a long-range Coulomb potential in addition to the usual Debye-screened form. With such an effective potential, we had investigated the effects of perturbative as well as nonperturbative effects in QGP on the dissociation of different quarkonium states [50, 51].

Now let us consider a central collision in a nucleus-nucleus collision, which results in formation of QGP at initial time  $\tau_i$ . We assume the plasma to cool, according to Bjorken's boost invariant longitudinal

expansion. There are many attempts [52-54] to incorporate the plasma expansion dynamics to study charmonium suppression in relativistic nuclear collisions. However, some important points were not addressed in their works [52-54] to quantify the suppression in an expanding system: the viscous forces in the energy-momentum tensor and the proper EoS. The effects of dissipative forces were not included in hydrodynamic expansion which cause the plasma to expand slowly and results in more suppression. The equation of state employed in their works was either ideal or bag model EoS used to calculate two vital factors - screening energy density,  $\epsilon_s$  of the system (corresponding to dissociation temperature) and the time,  $\tau_s$  elapsed by the system to reach at  $\epsilon_s$ , through expansion. Since the matter formed at RHIC is far from its ideal limit even at  $T \geq T_c$  so the ideal or bag model equation of state is not reliable to study the suppression of charmonium yields in an expanding strongly interacting system.

Recently we [55] had studied the survival of charmonium states in a dissipative strongly interacting QGP where the suppression of  $J/\psi$ 's due to potential screening in the deconfined medium at relativistic nuclear collisions is two step learning. First one is the understanding of dissociation in a static thermal medium for which knowledge of medium dependence of heavy quark potential is very much needed. Second one is the generalization of dissociation in an expanding medium where a equation of state plays the major role as an input. In [55], we used the equation of state for QGP, in analogy with strongly-coupled plasma where hadrons exist for  $T < T_c$  and go to plasma of quarks and gluons (QGP) for  $T > T_c$  and there is no hadrons or glue balls because confinement interactions due to QCD vacuum was assumed to be melted [1], although Zahed and Shuryak [31] suggested that QGP at temperatures up to few  $T_c$  supports weakly bound mesonic states. So the only interaction present in the deconfined plasma phase is Coulomb interaction and hence plasma parameter ( $\Gamma$ ) becomes the ratio of average Coulomb potential energy ( $=\frac{4}{3} \frac{\alpha_s}{r_{av}}$ ) to average kinetic energy ( $\sim T$ ). Finally, an expression for equation of state was obtained as a function of plasma parameter,  $\Gamma$  [1].

Using the abovementioned equation of state, we [55] got a better agreement with the PHENIX experimental results [56] compared to earlier works[53] but still it lacks complete agreement. The reason of disagreement may be two fold. This could be either due to the arbitrariness in the definition of dissociation criteria or due to the equation of state which is not fully compatible with the nonperturbative nature of QGP. In the present work, we revisited the abovementioned equation of state [1] by incorporating nonperturbative effects in the plasma parameter around which the equation of state is expanded.

As discussed above, the existence of nonperturbative interactions even at  $T \geq T_c$  indicates that the string tension may not vanish abruptly at  $T_c$ ,

so potential in the deconfined phase could have a nonvanishing confining (string) term, in addition to the Coulomb term [50] unlike Coulomb interaction alone in the aforesaid model [1]. This is the central theme of our work where EoS has been obtained by retaining both terms in the potential and then calculate the thermodynamic variables *viz* pressure, energy density, speed of sound etc. Our results match nicely with the lattice results of gluon [7], 2-flavor (massless) as well as 3-flavor (massless) QGP [57]. There is also an agreement with (2+1) (two massless and one is massive) and 4 flavoured lattice results too. Motivated by the agreement with lattice results, we employ our equation of state to study the  $J/\psi$  suppression in an expanding plasma in the presence of viscous forces with the universal ratio  $\eta/s = 1/4\pi$  [6]. We have found complete agreement with the PHENIX results on the centrality dependence of  $J/\psi$  suppression at RHIC energy. Further we have predicted the  $\Upsilon$  suppression which could be the potential candidate for the LHC experiment where  $J/\psi$  suppression could be marred by enhancement due to recombination of abundant  $c\bar{c}$  pairs.

The paper is organized as follows. In Sec.II.A, we briefly discuss our recent work on medium modified potential. Using this effective potential, we have then developed the equation of state for strongly interacting matter and have shown our results on pressure, energy density and speed of sound etc. along with the lattice data in Sec.II.B. In Sec.III.A, we have employed the aforesaid equation of state to study boost-invariant (1+1) dimensional longitudinal expansion in the presence of viscous forces and estimate the survival probability in a longitudinally expanding QGP in Sec.III.B. Results and discussion will be presented in Sec.IV and finally, we conclude in Sec.V.

## II. EQUATION OF STATE FOR STRONGLY-INTERACTING QGP

In this section, first we briefly discuss the medium modified effective potential in Sec.II.A. In Sec.II.B, we revisit the EoS of strongly coupled QGP developed by Bannur [1] and then obtain our EoS as a function of plasma parameter obtained from the medium-modified potential, discussed in Sec.II.A.

### A. Medium modified effective potential

In thermodynamical studies of QCD plasma phase, deviations from perturbative calculations and ideal gas behavior at temperatures much larger than the deconfinement temperature calls for quantitative non-perturbative calculations [45]. In light of this finding, one can not simply ignore the effects of string tension between the quark-antiquark pairs beyond  $T_c$ . Recently, this issue had successfully

been addressed in the context of the dissociation of quarkonium in QGP [50, 51] where we assumed the potential between a heavy quark-antiquark at  $T = 0$  as the Cornell potential :  $V(r) = -\alpha/r + \sigma r$  and then corrected its Fourier transform (FT)  $\tilde{V}(k)$ , to incorporate the medium modifications, as

$$\tilde{V}(k) = \frac{V(k)}{\epsilon(k)} \quad , \quad (1)$$

where  $\epsilon(k)$  is the dielectric permittivity [58], given by

$$\epsilon(k) = \left(1 + \frac{\Pi_L(0, k, T)}{k^2}\right) \equiv \left(1 + \frac{m_D^2}{k^2}\right) \quad , \quad (2)$$

and  $V(k)$  is the Fourier transform (FT) of the Cornell potential:

$$V(k) = -\sqrt{(2/\pi)} \frac{\alpha}{k^2} - \frac{4\sigma}{\sqrt{2\pi}k^4}. \quad (3)$$

Substituting Eqs.(2) and (3) into (1) and then evaluating its inverse FT one obtains the  $r$ -dependence of the medium modified potential [50, 51],

$$V(r) = \left(\frac{2\sigma}{m_D^2} - \alpha\right) \frac{e^{-m_D r}}{r} - \frac{2\sigma}{m_D^2 r} + \frac{2\sigma}{m_D} - \alpha m_D \quad , \quad (4)$$

where constant terms arise from the basic computations of real-time static potential [59] or from real- and imaginary-time correlators in a thermal QCD medium[60] and are introduced to yield the correct limit of  $V(r, T)$  as  $T \rightarrow 0$ . The potential (4) thus obtained has an additional long range Coulomb term with an (reduced) effective charge in addition to the conventional Yukawa term. It is worth to mention [61] that one-dimensional Fourier transform of the Cornell potential in the medium yields the screened form as used in the lattice QCD to study the quarkonium properties which assumes the one-dimensional color flux tube structure. However, at finite temperature that may not be the case since the flux tube structure may expand in more dimensions [11]. Therefore, it is better to consider the three-dimensional form of the medium modified Cornell potential.

Recently we had employed the above medium-modified potential to estimate the dissociation pattern of the charmonium and bottomonium states and also explore how the pattern changes as we go from perturbative to nonperturbative regime [50]. The results obtained in perturbative to nonperturbative domain were closer to the results obtained from the study of spectral function constructed in lattice QCD and potential model studies, respectively. This medium modified potential will be employed to develop the equation of state in next section.

## B. Equation of state

The equation of state for the quark matter produced in relativistic nucleus-nucleus collisions is an important observable and the properties of the matter are sensitive to it. The expansion of QGP is quite sensitive to EoS through the speed of sound which, in turn, explores the sensitivity of the quarkonium suppression to the equation of state [53, 54].

Lattice results [7] show that bulk thermodynamic observables such as pressure and energy density deviate from ideal gas behavior even at  $\sim 4 T_c$  and approach the ideal limit very slowly. This suggests that there are still (nonperturbative) interactions present in the deconfined medium. There have been many attempts to explain this strongly interacting matter formed during ultrarelativistic collisions. Many models, such as bag model [15], quasiparticle model [62, 63], strongly interacting quark-gluon plasma (QGP) model [14] were used to explain the complicated matter formed. None of the models were fully accepted by the physics community even though some models were successful to a great extent. An equation of state for such strongly interacting matter was developed by treating QGP near and above  $T_c$  in Cornell potential as Coulomb plus linear confinement term and compared with the lattice results for pure gauge, two-flavor, and three-flavor QGP [57]. Arnold and Zhai [17] and Zhai and Kastening [18] have developed a perturbative EoS of interacting quarks and gluons upto  $O(g^5)$  which has further been extended by Kajantie et al. [19, 20] by including the nonperturbative contributions *viz.*  $O(gT)$  and  $O(g^2T)$  up to  $O(g^6 \ln(1/g))$  [21]. On the other hand, a semiclassical approach is devised to study the bulk properties of QGP [22, 23, 64–66] where nonperturbative effects manifested as effective mean color fields. These color fields have the dual role of producing the soft and semisoft partons, apart from modulating their interactions. The emergence of such effective field degrees of freedom, together with a classical transport has been indicated earlier [64, 65].

Recently Bannur [1] developed an equation of state for a strongly-coupled QGP by appropriate modifications of strongly-coupled plasma in QED to take account color and flavor degrees of freedom with the running coupling constant and a reasonably good fit to the lattice results was obtained. Let us briefly discuss the strongly coupled plasma in QED where the equation of state is expressed as a function of plasma parameter  $\Gamma$  [67]:

$$\epsilon_{\text{QED}} = \left( \frac{3}{2} + u_{ex}(\Gamma) \right) n T, \quad (5)$$

where the first term represents the ideal contribution and the deviations from ideal EoS is given by,

$$u_{ex}(\Gamma) = \frac{u_{ex}^{Abe}(\Gamma) + 3 \times 10^3 \Gamma^{5.7} u_{ex}^{OCP}(\Gamma)}{1 + 3 \times 10^3 \Gamma^{5.7}}, \quad (6)$$

where  $u_{ex}^{Abe}$  is given by

$$u_{ex}^{Abe}(\Gamma) = -\frac{\sqrt{3}}{2} \Gamma^{3/2} - 3 \Gamma^3 \left[ \frac{3}{8} \ln(3\Gamma) + \frac{\gamma}{2} - \frac{1}{3} \right], \quad (7)$$

which was derived by Abe [68] in the formalism of giant cluster expansion with the Euler constant  $\gamma$  and is valid for  $\Gamma < .1$ . The term  $u_{ex}^{OCP}$  is given by

$$u_{ex}^{OCP} = -0.898004\Gamma + 0.96786\Gamma^{1/4} + 0.220703\Gamma^{-1/4} - 0.86097, \quad (8)$$

which was numerically obtained for one component plasma and is valid all  $\Gamma < 180$  [69]. In the above equations, the plasma parameter  $\Gamma$  is the ratio of average potential energy to the average kinetic energy.

Let us now consider strongly-coupled plasma in QCD where it was assumed that hadron exists for  $T < T_c$  and goes to QGP for  $T > T_c$ . That is, for  $T > T_c$ , it is the strongly interacting plasma of quarks and gluons and no hadrons or glue balls because it was assumed that confinement interactions due to QCD vacuum has been melted [1] at  $T = T_c$ . Hence the only interaction present in the deconfined plasma phase is the Coulomb interaction and so the plasma parameter ( $\Gamma$ ) was evaluated as the ratio of average Coulomb potential energy to average kinetic energy. As discussed earlier in Sec.II.A, we will retain confinement interactions, in addition to the Coulomb interaction, through the linear term in the potential which manifests in the nonzero values of string tension even at  $T \geq T_c$  in the potential (4). Finally, the equation of state has been obtained by using the potential (4) in the plasma parameter after inclusion of relativistic and quantum effects as:

$$\epsilon = \left( 3 + u_{ex}(\Gamma) \right) n T, \quad (9)$$

where the form of  $u_{ex}(\Gamma)$  remains the same as in (6). In terms of ideal contribution, the scaled-energy density is written as

$$e(\Gamma) \equiv \frac{\epsilon}{\epsilon_{SB}} = 1 + \frac{1}{3} u_{ex}(\Gamma), \quad (10)$$

where  $\epsilon_{SB}$  is given by,

$$\epsilon_{SB} \equiv 3a_f T^4, \quad (11)$$

where  $a_f \equiv (16 + 21 n_f/2)\pi^2/90$  is a constant which depends on degrees of freedom of quarks and gluons. Here we have employed the QCD running coupling in  $\overline{\text{MS}}$  scheme [70], in compatible with lattice simulation, up to two-loop level:

$$g^2 \approx 2b_0 \ln \frac{\bar{\mu}}{\Lambda_{\overline{\text{MS}}}} \left( 1 + \frac{b_1}{2b_0^2} \frac{\ln \left( 2 \ln \frac{\bar{\mu}}{\Lambda_{\overline{\text{MS}}}} \right)}{\ln \frac{\bar{\mu}}{\Lambda_{\overline{\text{MS}}}}} \right)^{-1}, \quad (12)$$



where  $b_0 = (33 - 2n_f)/(48\pi^2)$  and  $b_1 = (153 - 19n_f)/(384\pi^4)$ ,  $\bar{\mu}$  and  $\Lambda_{\overline{\text{MS}}}$  are the scale parameter and the renormalization scale in  $\overline{\text{MS}}$  scheme, respectively. For, the EoS to depend on the renormalization scale, the physical observables should be scale independent. We circumvent the problem by trading off the dependence on  $\Lambda_{\overline{\text{MS}}}$  to a dependence on the critical temperature  $T_c$ .

$$\begin{aligned} \bar{\mu} \exp(\gamma_E + c) &= \Lambda_{\overline{\text{MS}}}(T) \\ \Lambda_{\overline{\text{MS}}}(T) \exp(\gamma_E + c) &= 4\pi\Lambda_T \quad , \end{aligned} \quad (13)$$

where  $c$  is a constant depending on colors and flavors:  $c = (n_c - 4n_f \ln 4) / (22n_c - n_f)$  and  $\gamma_E = 0.5772156$ . There are several ambiguities, associated with the renormalization scale  $\Lambda_{\overline{\text{MS}}}$ , the scale parameter  $\bar{\mu}$  which occurs in the expression for the running coupling constant  $\alpha_s$ . This issue has been discussed well in literature and a popular way out is the BLM criterion due to Brodsky, Lepage and Mackenzie [71]. In this criterion,  $\Lambda_{\overline{\text{MS}}}$  is allowed to vary between  $\pi T$  and  $4\pi T$  [72]. For our purposes, we choose the renormalization scale  $\Lambda_{\overline{\text{MS}}}$  close to the central value  $2\pi T_c$  [73] for  $n_f=0$  and  $\pi T_c$  for both  $n_f=2$  and  $n_f=3$  flavors.

It is worth to mention here that if the factor  $\frac{b_1}{2b_0^2} \frac{\ln(2 \ln \frac{\bar{\mu}}{\Lambda_{\overline{\text{MS}}}})}{\ln \frac{\bar{\mu}}{\Lambda_{\overline{\text{MS}}}}}$  is much smaller than 1 ( $\ll 1$ ) then the above expression reduces to the expression used in [1, Eq.(10)], after neglecting the higher order terms of the above factor. However, this possibility does not hold good for the temperature ranges used in the calculation and cause an error in coupling which finally makes the difference in the results between our model and Bannur model [1]. Finally, we get the energy density  $\varepsilon(T)$  from Eq.(10) and using the thermodynamic relation,

$$\varepsilon = T \frac{dp}{dT} - P \quad , \quad (14)$$

we get the pressure as

$$\frac{P}{T^4} = \left( \frac{P_0}{T_0} + 3a_f \int_{T_0}^T d\tau \tau^2 e(\Gamma(\tau)) \right) / T^3 \quad , \quad (15)$$

where  $P_0$  is the pressure at some reference temperature  $T_0$  and has been fixed with the values of pressure at critical temperature for a particular system - gluon plasma, 2-flavor plasma etc. Once we know the pressure  $P$  and energy density  $\varepsilon$ , the speed of sound  $c_s^2 (= \frac{dP}{d\varepsilon})$  can be evaluated.

In Fig. 1, we have plotted the variation of pressure ( $P/T^4$ ) with temperature ( $T/T_c$ ) for pure gauge, 2-flavor and 3-flavor QGP along with lattice results. In Bannur model [1], for each system,  $g_c$  and  $\Lambda_T$  are adjusted to get a good fit to lattice results. However, in our calculation, there is no quantity to be fitted

for predicting lattice results. We have fixed  $P_0$  from the lattice data at the critical temperature  $T_c$  for each system, separately.

Once pressure,  $P(T)$  is obtained, then other macroscopic quantities such as energy density  $\varepsilon$ , speed of sound  $c_s^2$  etc. can be derived from  $P(T)$  and no other parameters are needed. In Fig. 2, we plotted the energy density ( $\varepsilon/T^4$ ) with temperature ( $T/T_c$ ) for all three systems along with lattice results and a reasonably good fit is obtained without any extra parameters. All three curves look similar, but shifts to left as flavor content increases. For the sake of comparison with the results of Bannur EoS, we have taken the critical temperatures  $T_c$  equal to 275, 175 and 155 MeV for gluon plasma, 2-flavor and 3-flavor QGP, respectively.

In Fig. 3, the speed of sound,  $c_s^2$  is plotted for all three systems, but matching have been checked with the only available lattice results for gluon plasma. There is an excellent agreement with the lattice results and we have predicted the results for the flavored QGP. All three curves have similar behaviour, i.e, sharp rise near  $T_c$  and then flatten to the ideal value (1/3). However,  $c_s^2$  is larger for larger flavor content of plasma. In the vicinity of critical temperature, fits or predictions may not be good, especially for energy density  $\varepsilon$  and  $c_s^2$  which strongly depends on variations of pressure  $P$  with respect to temperature  $T$ . Lattice data also has large error bars very close to  $T_c$ . However, except for small region at  $T = T_c$ , our results are very good for all regions of  $T > T_c$ . It is interesting to note that recently Peshier and Cassing [74] also obtained similar results on the dependence of plasma parameter  $\Gamma$  in quasi-particle model and concluded that QGP behaves like a liquid, not weakly-interacting gas.

Now we are interested to study for the realistic case where u and d quarks have very small masses (5-10 MeV) and strange quarks are having masses 150-200 MeV and charm quark with mass 1.5 GeV. Let  $g_f$  counts the effective number of degrees of freedom of a massive Fermi gas. For a massless gas we have, of course,  $g_f = n_f$ . In general we define

$$n_f = \sum_{f=u,d,\dots} g(m_f/T) \quad (16)$$

where,

$$g\left(\frac{m_f}{T}\right) = \frac{360}{7\pi^4} \int_0^\infty dx \, x \sqrt{x^2 - \left(\frac{m_f}{T}\right)^2} \times \ln(1 + e^{-x}) \quad (17)$$

In Fig.4, we have shown our results on (2+1)-flavors and 4-flavors QGP along with lattice data [75, 76] and replotted the variation of  $P(T)/T^4$  with temperature  $T/T_c$  for all systems. Similar plots for energy density  $\varepsilon(T)/T^4$  with temperature  $T/T_c$  for all systems is replotted in Fig. 5. We have also compared with the results from Bannur model.

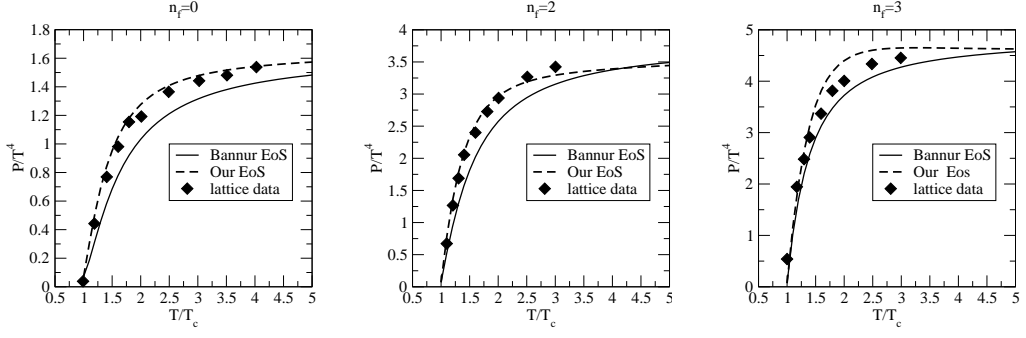


FIG. 1: Plots of  $P/T^4$  as a function of  $T/T_c$  from Bannur EoS, our EoS and lattice results for (a) pure gauge (extreme left figure), 2-flavor QGP (middle figure) and 3-flavor QGP (extreme right figure). In each figure, solid line represents the results obtained from Bannur EoS, dashed line represents the results from our EoS and diamond symbols represent lattice results.

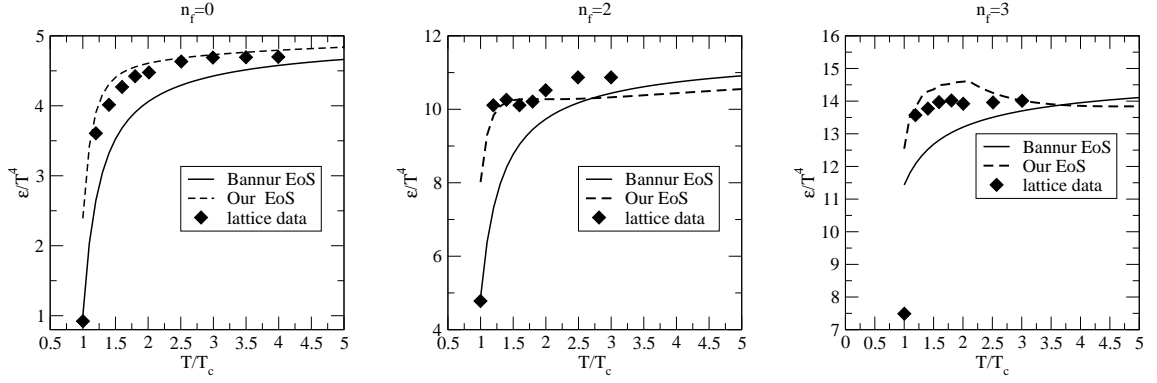


FIG. 2: Plots of  $\epsilon/T^4$  as a function of  $T/T_c$ . The notations are the same as in Fig.1.

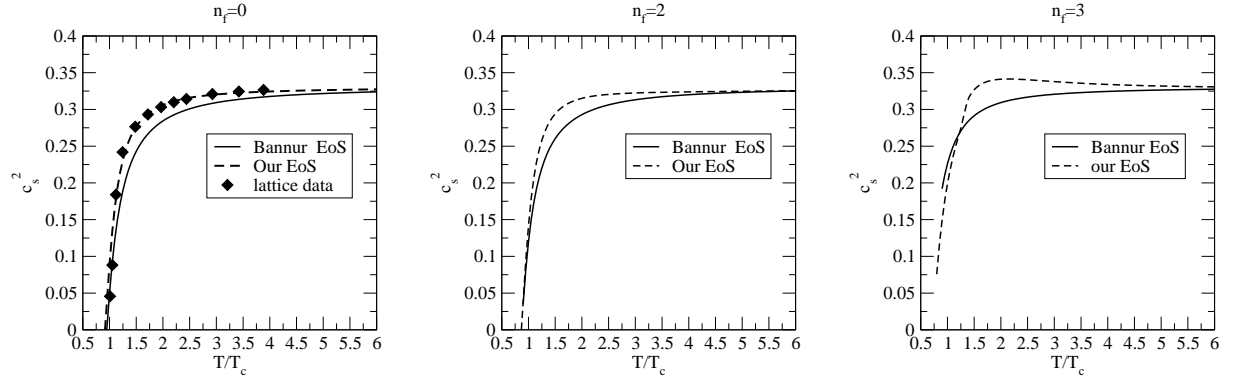


FIG. 3: Plots of  $c_s^2$  as a function of  $T/T_c$  from our EoS and Bannur EoS for pure gauge, 2-flavor QGP and 3-flavor QGP.

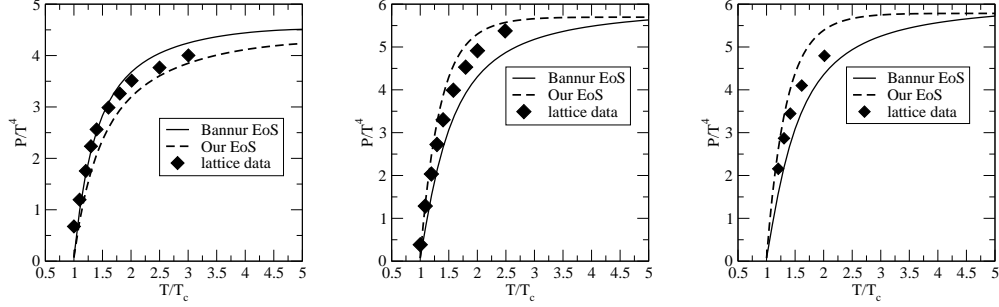


FIG. 4: variation of  $P/T^4$  as a function of  $T/T_c$  for a) two massless and one massive (2+1), b) and c) for 4-flavour QGP for two different masses,  $m/T=0.4$  and  $0.2$ , respectively. The notations are the same as in Fig.1.

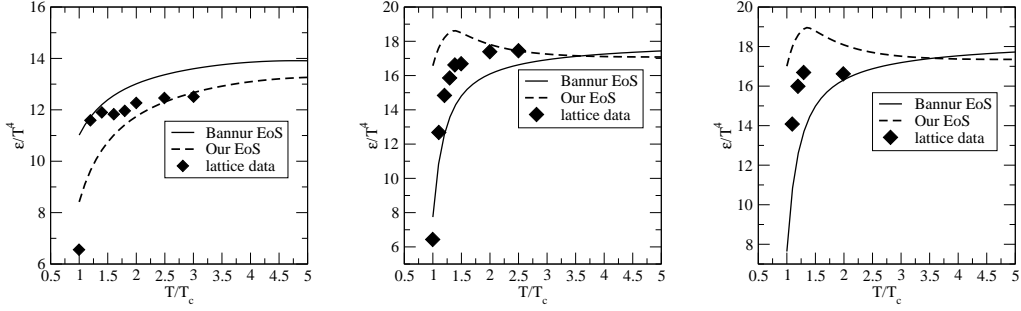


FIG. 5: Variation of  $\epsilon/T^4$  as a function of  $T/T_c$  where the notations are the same as in Fig.4.

This indicates that in the presence of a heavier quark the deviations of the pressure from the ideal gas value is larger than in the massless limit. This is in qualitative agreement with the observations. In view of the agreement with the results of lattice equation of state, our EoS is a right choice for the strongly-interacting matter possibly formed at RHIC to calculate the thermodynamical quantities *viz.* screening energy density ( $\epsilon_s$ ), the speed of sound etc. to study the hydrodynamical expansion of plasma and finally, to estimate the suppression of  $J/\psi$  in nuclear collisions.

### III. SUPPRESSION OF $J/\psi$ IN A LONGITUDINALLY EXPANDING PLASMA

In Sec.III.A, we study hydrodynamic boost-invariant Bjorken expansion in (1+1) dimension with the EoS discussed in Sec.IIB as an input. In addition, we explore the effects of dissipative terms up to first-order in the stress-tensor. Then we turn our attention to derive the  $J/\psi$  survival probability for an expanding QGP, in Sec.III.B.

#### A. Longitudinal expansion in the presence of dissipative forces

In the presence of viscous forces, the energy-momentum tensor is written as,

$$T^{\mu\nu} = (\epsilon + p)u^\mu u^\nu + g^{\mu\nu}p + \pi^{\mu\nu}, \quad (18)$$

where the stress-energy tensor,  $\pi^{\mu\nu}$  up to first-order is given by

$$\pi^{\mu\nu} = \eta \langle \nabla^\mu u^\nu \rangle, \quad (19)$$

where  $\eta$  is the co-efficient of the shear viscosity and  $\langle \nabla^\mu u^\nu \rangle$  is the symmetrized velocity gradient.

In (1+1) dimensional Bjorken expansion in the first-order dissipative hydrodynamics, only one component  $\pi^{\eta\eta}$  of the viscous stress tensor is non-zero, hence the equation of motion reads,

$$\partial_\tau \epsilon = -\frac{\epsilon + p}{\tau} + \frac{4\eta}{3\tau^2}. \quad (20)$$

The first term in the RHS is the same as in the case of zeroth-order (non-viscous) hydrodynamics and the second term is the correction arising from constant  $\eta/s$  which causes the system to expand slowly compared to the perfect fluid,  $\eta = 0$ .

The solution of equation of motion (19) is obtained as,

$$\begin{aligned} \epsilon(\tau)\tau^{(1+c_s^2)} + \frac{4a}{3\tilde{\tau}^2}\tau^{(1+c_s^2)} &= \epsilon(\tau_i)\tau_i^{(1+c_s^2)} + \frac{4a}{3\tilde{\tau}_i^2} \\ &= \text{const}, \end{aligned} \quad (21)$$

where the constant  $a$  is  $(\frac{2}{s})T_i^3\tau_i$  and the symbols,  $\tilde{\tau}^2$  and  $\tilde{\tau}_i^2$  are given by  $(1-c_s^2)\tau^2$  and  $(1-c_s^2)\tau_i^2$ , respectively. The first term accounts for the contributions coming from the zeroth-order expansion (ideal fluid) and the second term is the first-order viscous corrections. The above equation (21) will finally be coupled with the dissociation of a quarkonium in a static thermal medium to estimate the quarkonium suppression in relativistic nucleus-nucleus collisions.

#### B. Survival probability

We now have all the ingredients to write down the survival probability. Chu and Matsui [52] studied the transverse momentum dependence ( $p_T$ ) of the



survival probability by choosing the speed of sound  $c_s^2 = 1/3$  (ideal EoS) and the extreme value  $c_s^2 = 0$ . This work was further generalized by invoking the various parameters for Au-Au collisions at RHIC in [53] to include the effects of realistic EoS in an adhoc manner by simply choosing a lower value of  $c_s^2 = 1/5$ . Instead of taking arbitrary values of  $c_s^2$  we tabulated the values of  $c_s^2$  in Tables II and IV-VI corresponding to the dissociation temperatures calculated from our EoS [55] for charmonium and bottomonium states, respectively. Moreover, in the light of recent experimental finding from RHIC, one cannot ignore the viscous effects while studying charmonium suppression. Here, we address these issues.

Let us take an initial energy-density profile on a transverse plane :

$$\epsilon(r; \tau_i) = \epsilon_i \left(1 - \frac{r^2}{R_T^2}\right)^\beta \Theta(R_T - r) \quad (22)$$

where  $r$  is the transverse co-ordinate and  $R_T$  is the transverse radius of the nucleus. One can define an average energy density  $\langle \epsilon_i \rangle$  as

$$\pi R_T^2 \langle \epsilon_i \rangle = \int 2\pi r dr \epsilon(r; \tau_i) \quad (23)$$

so that

$$\epsilon_i = (1 + \beta) \langle \epsilon_i \rangle \quad ; \beta = 1. \quad (24)$$

The average initial energy density  $\langle \epsilon_i \rangle$  [79] will be given by the modified Bjorken formula:

$$\langle \epsilon_i \rangle = \frac{\xi}{A_T \tau_i} \left( \frac{dE_T}{dy_h} \right)_{y_h=0}, \quad (25)$$

where  $A_T$  is the transverse overlap area of the colliding nuclei and  $(dE_T/dy_h)_{y_h=0}$  is the transverse energy deposited per unit rapidity of outgoing hadrons. Both depend on the number of participants  $N_{part}$  [80] and thus provide centrality dependent initial average energy density  $\langle \epsilon_i \rangle$  in the transverse plane (Table I). For this purpose, we have extracted the transverse overlap area  $A_T$  and the pseudo-rapidity distribution  $dE_T/d\eta_h|_{\eta_h=0}$  [80] at various values of number of participants  $N_{part}$ . These  $dE_T/d\eta_h|_{\eta_h=0}$  numbers are then multiplied by a Jacobian 1.25 to yield the rapidity distribution  $dE_T/dy_h|_{y_h=0}$  which will be further used to calculate the average initial energy density from Bjorken formula (25). The scaling factor  $\xi = 5$  has been introduced in order to obtain the desired values of initial energy densities [81] for most central collision which are consistent with the predictions of the self-screened parton cascade model [82] and also with the requirements of hydrodynamic simulation [81] to fit the pseudo-rapidity distribution of charged particle multiplicity  $dN_{ch}/d\eta$  for various centralities

TABLE I: Kinematic characterization of Au+Au collisions at RHIC [56]

Nuclei	$\sqrt{s_{NN}}$ (GeV)	$N_{part}$	$\langle \epsilon_i \rangle$ (GeV/fm <sup>3</sup> )	$R_T$ (fm)
Au+Au 200		22.0	5.86	3.45
		30.2	7.92	3.61
		40.2	10.14	3.79
		52.5	12.76	3.96
		66.7	15.69	4.16
		83.3	18.58	4.37
		103.0	21.36	4.61
		125.0	24.38	4.85
		151.0	27.37	5.12
		181.0	30.52	5.38
		215.0	34.17	5.64
		254.0	37.39	5.97
		300.0	41.08	6.31
		353.0	45.09	6.68

observed in PHENIX experiments at RHIC energy. Later we will discuss the centrality dependence of initial energy densities at LHC energy in Sec.IV.

The (screening) time,  $\tau_s$  when the energy density of the system drops to the screening energy density  $\epsilon_s$  is estimated from Eq.(21) as

$$\tau_s(r) = \tau_i \left[ \frac{\epsilon_i(r) - \frac{4a}{3\tilde{\tau}_i^2}}{\epsilon_s - \frac{4a}{3\tilde{\tau}_s^2}} \right]^{\frac{1}{1+c_s^2}} \quad (26)$$

where  $\epsilon_i(r) \equiv \epsilon(\tau_i; r)$  and  $\tilde{\tau}_s^2$  is  $(1 - c_s^2)\tau_s^2$ . The critical radius  $r_s$ , is seen to mark the boundary of the region where the quarkonium formation is suppressed, can be obtained by equating the duration of screening  $\tau_s(r)$  to the formation time  $t_F = \gamma\tau_F$  for the quarkonium in the plasma frame and is given by:

$$r_s = R_T(1 - A)^{1/2} \Theta(1 - A), \quad (27)$$

where  $A$  is given by

$$A = \left[ \left( \frac{\epsilon_s}{\epsilon_i} \right) \left( \frac{t_F}{\tau_i} \right)^{1+c_s^2} + \frac{1}{\epsilon_i} \left( \frac{t_F}{\tau_i} \right)^{(1+c_s^2)} \frac{4a}{3\tilde{t}_F^2} + \frac{1}{\epsilon_i} \frac{4a}{3\tilde{\tau}_i^2} \right]^{1/\beta} \quad (28)$$

with  $\tilde{t}_F^2 = (1 - c_s^2)t_F^2$ . The quark-pair will escape the screening region and form quarkonium if its position vector  $\mathbf{r}$  and transverse momentum  $\mathbf{p}_T$  are such that

$$|\mathbf{r} + \tau_F \mathbf{p}_T / M| \geq r_s. \quad (29)$$

Thus, if  $\phi$  is the angle between the vectors  $\mathbf{r}$  and  $\mathbf{p}_T$ , then the above condition reduces to

$$\cos \phi \geq [(r_s^2 - r^2) M - \tau_F^2 p_T^2 / M] / [2 r \tau_F p_T], \quad (30)$$

which leads to a range of values of  $\phi$  when the quarkonium would escape. Now we can write for the survival probability of the quarkonium:

$$S(p_T) = \frac{\left[ \int_0^{R_T} r dr \int_{-\phi_{\max}}^{+\phi_{\max}} d\phi P(\mathbf{r}, \mathbf{p}_T) \right]}{\left[ 2\pi \int_0^{R_T} r dr P(\mathbf{r}, \mathbf{p}_T) \right]}, \quad (31)$$

where  $\phi_{\max}$  is the maximum positive angle ( $0 \leq \phi \leq \pi$ ) allowed by Eq.(30):

$$\phi_{\max} = \begin{cases} \pi & \text{if } y \leq -1 \\ \cos^{-1} |y| & \text{if } -1 < y < 1 \\ 0 & \text{if } y \geq 1 \end{cases}, \quad (32)$$

where

$$y = [(r_s^2 - r^2) M - \tau_F^2 p_T^2 / M] / [2 r \tau_F p_T], \quad (33)$$

and  $P$  is the probability for the quark-pair production at  $(\mathbf{r}, \mathbf{p}_T)$ , in a hard collision which may be factored out as

$$P(\mathbf{r}, \mathbf{p}_T) = f(r)g(p_T), \quad (34)$$

where we take the profile function  $f(r)$  as

$$f(r) \propto \left[ 1 - \frac{r^2}{R_T^2} \right]^\alpha \Theta(R_T - r) \quad (35)$$

with  $\alpha = 1/2$ .

Often experimental measurement of survival probability at a given number of participants ( $N_{\text{part}}$ ) or rapidity ( $y$ ) is reported in terms of the  $p_T$ -integrated yield ratio (nuclear modification factor) over the range  $p_T^{\min} \leq p_T \leq p_T^{\max}$  whose theoretical expression would be

$$\langle S(p_T) \rangle = \frac{\int_{p_T^{\min}}^{p_T^{\max}} dp_T S(p_T)}{\int_{p_T^{\min}}^{p_T^{\max}} dp_T} \quad (36)$$

In nucleus-nucleus collisions, it is known that only about 60% of the observed  $J/\psi$  originate directly in hard collisions while 30% of them come from the decay of  $\chi_c$  and 10% from the decay of  $\psi'$ . Hence, the  $p_T$ -integrated inclusive survival probability of  $J/\psi$  in the QGP becomes [11, 53].

$$\langle S^{\text{incl}} \rangle = 0.6 \langle S^{\text{dir}} \rangle_\psi + 0.3 \langle S^{\text{dir}} \rangle_{\chi_c} + 0.1 \langle S^{\text{dir}} \rangle_{\psi'} \quad (37)$$

#### IV. RESULTS AND DISCUSSIONS

Before displaying the results, let us discuss the physical understanding of quarkonium suppression due to screening in the deconfined medium produced in relativistic nucleus-nucleus collisions. This involves a competition of various time-scales involved in an expanding plasma. First one is the screening time,  $\tau_s$  as the time available for the hot and dense expanding system during which  $J/\psi$ 's are suppressed. Second one is the formation time of  $J/\psi$  in the plasma frame ( $t_F = \gamma \tau_F$ ) which depends on the transverse momentum by which the  $c\bar{c}$  pairs was originally produced. Third one is the cooling rate which depends on the speed of sound,  $c_s^2$  through the equation of state. The screening time not only depends upon the screening energy density,  $\epsilon_s$  but also depends on the speed of sound through equation of state. The value of  $\epsilon_s$  is different for different charmonium states and is calculated from the equation of state and hence varies from one EoS to other. If  $\epsilon_s \gtrsim \epsilon_i$ , initial energy density, then there will be no suppression at all i.e., survival probability,  $S(p_T)$  is equal to 1.

More precisely, the screening time depends upon i) the screening energy density and the difference between a given initial energy density  $\epsilon_i$  and screening energy density  $\epsilon_s$ : the more will be the difference the more will be the suppression, ii) the speed of sound: for  $c_s^2$  less than the ideal limit (1/3), the rate of cooling will be slower which, in turn, makes the screening time larger for a fixed difference in ( $\epsilon_i - \epsilon_s$ ) and leads to more suppression, and iii) the  $\eta/s$  ratio: an additional handle to explore the equation of state by controlling the expansion of the plasma. If the ratio is non-zero then the cooling will be slower compared to  $\eta/s=0$ , so the system will take longer time to reach  $\epsilon_s$  resulting the higher value of screening time and hence more suppression compared to  $\eta/s = 0$ . With this physical understanding we analyze our results,  $\langle S(p_T) \rangle$  as a function of the number of participants  $N_{\text{part}}$  in an expanding QGP.

In our analysis, we have employed the dissociation temperatures for the charmonium states ( $J/\psi$ ,  $\chi_c$  etc.) computed from the lattice QCD correlator studies [85] in Table II. The corresponding values of screening energy densities,  $\epsilon_s$  and the speed of sound  $c_s^2$  calculated in our EoS are also listed which will be used as inputs along with the kinematic data in Table I, to calculate  $\langle S(p_T) \rangle$ . We have shown the variation of  $p_T$ -integrated survival probability (in the range allowed by invariant  $p_T$  spectrum of  $J/\psi$  in Phenix experiment [56]) with  $N_{\text{part}}$  at mid-rapidity in Fig.6. The experimental data (the nuclear-modification factor  $R_{AA}$ ) are shown by the squares with error bars whereas circles represent sequential suppression. It is found that our calculation matches completely with the experimental results. To see the importance of confinement interactions in the deconfined phase, we have also calcu-

TABLE II: Formation time (fm), dissociation temperature  $T_D$  [83], the speed of sound  $c_s^2$  and the screening energy density  $\epsilon_s$  ( $\text{GeV}/\text{fm}^3$ ) for charmonium states, calculated both in our and Bannur EoS, respectively.

State	$\tau_F$	$T_D$	$c_s^2(\text{our})$	$c_s^2(\text{BAN})$	$\epsilon_s(\text{our})$	$\epsilon_s(\text{BAN})$
$J/\psi$	0.89	2.1	0.308	0.275	29.33	32.05
$\psi'$	1.50	1.12	0.255	0.214	01.94	02.36
$\chi_c$	2.00	1.16	0.261	0.220	02.28	0.220

lated the  $p_T$ -integrated survival probability (denoted by diamonds) using the equation of state in Bannur model [1] where there are no confinement interactions in the deconfined phase and have not found the agreement, as seen in our case. This difference is due to the fact that the screening energy density calculated in our EoS is smaller than the value obtained from Bannur model (as seen in Table II). The smaller value of screening energy density  $\epsilon_s$  causes an increase in the screening time and results in more suppression to match with the PHENIX results at RHIC.

At RHIC energy,  $J/\psi$  yields have been resulted from a balance between annihilation of  $J/\psi$ 's due to hard, thermal gluons [86, 87] along with colour screening [52, 88] and enhancement due to coalescence of uncorrelated  $c\bar{c}$  pairs [89–91] which are produced thermally at deconfined medium, at the phase boundary during (statistical) hadronization [92, 93]. However, recent PHENIX data do not show a fully confirmed indication of  $J/\psi$  enhancement except for the fact that  $\langle p_T^2 \rangle$  of the data and shape of rapidity-dependent nuclear modification factor  $R_{AA}(y)$  [56] show some characteristics of coalescence production.

On the other hand, at LHC energy, there will be an abundance of thermally produced  $c\bar{c}$  pairs so that one cannot make any concrete conclusion about the possible formation of QGP whereas the number of  $b\bar{b}$  pairs produced in hot deconfined medium will be meagre (because of large bottom quark-antiquark masses) so that the competition between the suppression due to screening and dissociation both and the enhancement due to recombination may be unlikely. Therefore, it will be interesting to estimate the  $\Upsilon$  suppression at LHC energy because the initial energy density at ALICE experiments will be high enough to suppress the  $\Upsilon$ 's sequentially making the bottomonium suppression an unambiguous signature.

Before finding the centrality (or impact parameter) dependence of  $\Upsilon$  suppression at LHC energy, it is necessary to know the centrality dependence of average initial energy density  $\langle \epsilon_i \rangle$  in terms of the number of participants  $N_{part}$  at LHC energy. Recently Eskola et al [82] computed the dependence of initial energy density, gluon, quark and antiquark

TABLE III: Kinematic characterization of Pb+Pb collisions at LHC [82]

Nuclei	$\sqrt{s_{NN}}$ (GeV)	$N_{part}$	$\langle \epsilon_i \rangle$ (GeV/fm <sup>3</sup> )	$R_T$ (fm)
Pb+Pb 5500	22.0		217.97	3.45
	30.2		236.12	3.61
	40.2		253.15	3.79
	52.5		269.11	3.96
	66.7		291.45	4.16
	83.3		315.55	4.37
	103.0		341.05	4.61
	125.0		370.49	4.85
	151.0		400.42	5.12
	181.0		433.18	5.38
	215.0		463.66	5.64
	254.0		507.04	5.97
	300.0		550.88	6.31
	353.0		600.39	6.68

numbers produced in ultrarelativistic heavy ion collisions with beam energy in parton saturation model. This model is based on the argument that the effects of all momentum scales can be estimated by performing the computation at the saturation momentum scales. The main emphasis of the study was at LHC and RHIC energies, although it gives reasonable good results at SPS too. The dependence of atomic number and beam energies of initial energy density, number density, temperature etc. are given by [82]

$$\epsilon_i = 0.103 \text{ GeV fm}^{-3} A^{0.504} (\sqrt{s})^{0.786}, \quad (38)$$

$$n_i = 0.370 \text{ fm}^{-3} A^{0.383} (\sqrt{s})^{0.574}, \quad (39)$$

$$T_i = 0.111 \text{ GeV } A^{0.126} (\sqrt{s})^{0.197} \quad (40)$$

Using this model, we have first calculated the centrality dependence of the initial energy densities at RHIC energy and found a good agreement with the known values (Table I) obtained from the centrality dependence of observed particle multiplicities [80]. Inspired by the success of the parton saturation model at RHIC energy, we have provided the centrality dependence of initial conditions at LHC energy in Table III and predict the suppression of  $\Upsilon$  yields at LHC energy, which is yet to be verified with the results available in ALICE experiments at LHC energy. We have taken three sets of dissociation temperatures for the bottomonium states [84]. These are obtained by equating a) the lattice free energy, b) by subtracting entropy term from the lattice free energy and c) linear combination of both a) and b), with the temperature dependent heavy quark effective potential. In Fig. 7, we have plot-

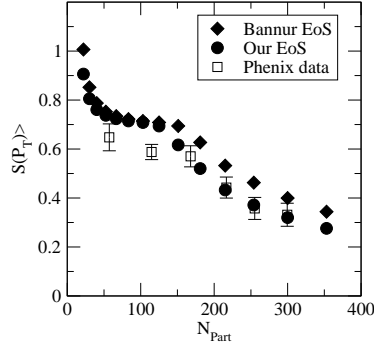


FIG. 6: The variation of  $p_T$  integrated survival probability (in the range allowed by invariant  $p_T$  spectrum of  $J/\psi$  by the Phenix experiment [56]) versus number of participants at mid-rapidity. The experimental data (the nuclear-modification factor  $R_{AA}$ ) are shown by the squares with error bars whereas circles and diamonds represent sequential melting using the values of  $T_D$ 's [50] and related parameters from Table II.

TABLE IV: Formation time (fm), dissociation temperature  $T_D$  [84], the speed of sound  $c_s^2$  and the screening energy density  $\epsilon_s$  ( $GeV/fm^3$ ) calculated in our EoS for bottomonium states, respectively.

State	$\tau_F$	$T_D$	$c_s^2(\text{our})$	$\epsilon_s(\text{our})$
$\Upsilon$	0.76	4.18	0.322	496.86
$\Upsilon'$	1.90	1.47	0.287	6.31
$\chi_b$	2.60	1.61	0.294	9.38

TABLE V: Formation time (fm), dissociation temperature  $T_D$  [84], the speed of sound  $c_s^2$  and the screening energy density  $\epsilon_s$  ( $GeV/fm^3$ ), calculated in our EoS for bottomonium states, respectively.

State	$\tau_F$	$T_D$	$c_s^2(\text{our})$	$\epsilon_s(\text{our})$
$\Upsilon$	0.76	3.40	0.320	213.14
$\Upsilon'$	1.90	1.18	0.263	2.44
$\chi_b$	2.60	1.22	0.267	2.82

ted the centrality dependence of the  $p_T$ -integrated survival probability for the bottomonium states at LHC energy.

TABLE VI: Formation time (fm), dissociation temperature  $T_D$  [84], the speed of sound  $c_s^2$  and the screening energy density  $\epsilon_s$  ( $GeV/fm^3$ ), calculated in SQGP EoS for bottomonium states, respectively.

State	$\tau_F$	$T_D$	$c_s^2(\text{our})$	$\epsilon_s(\text{our})$
$\Upsilon$	0.76	2.90	0.317	111.29
$\Upsilon'$	1.90	1.06	0.244	1.47
$\chi_b$	2.60	1.07	0.247	1.58

## V. CONCLUSIONS

We revisited the equation of state for strongly interacting quark-gluon plasma in the framework of strongly coupled plasma with appropriate modifications to take account of color and flavor degrees of freedom and QCD running coupling constant. In addition, we incorporate the nonperturbative effects in terms of nonzero string tension in the deconfined phase, unlike the Coulomb interactions alone in the deconfined phase beyond the critical temperature. Our results on thermodynamic observables *viz.* pressure, energy density, speed of sound etc. nicely fit the results of lattice equation of state with gluon, massless and as well *massive* flavored plasma. Motivated by this agreement we apply our equation of state to estimate the centrality dependence of  $J/\psi$  suppression in an expanding dissipative strongly in-

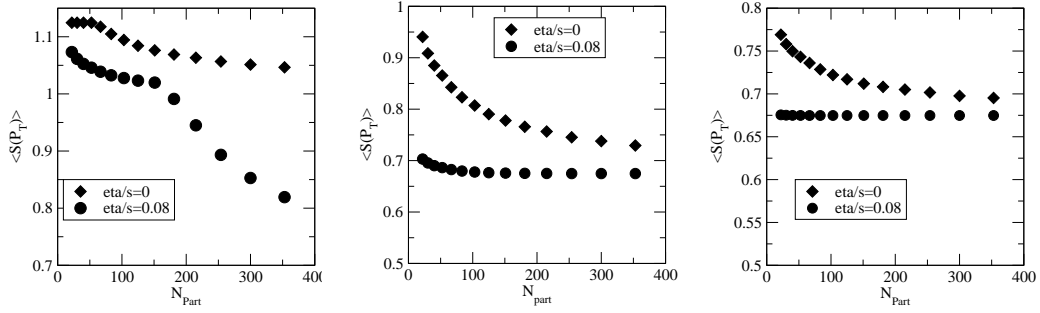


FIG. 7: The variation of  $p_T$  integrated survival probability versus number of participants for  $\Upsilon$ . The circles and diamonds represent sequential melting of  $\eta/s = 1/4\pi$  and  $\eta/s = 0$ , respectively. The parameter for left, middle and right figures given in the Table IV, V and VI, respectively.

interacting QGP produced in relativistic heavy-ion collisions. We have found a complete agreement with the PHENIX experimental results on  $J/\Psi$  suppres-

sion at RHIC energy. Moreover we predicted the same for the  $\Upsilon$  suppression which is yet to be verified in ALICE experiments at LHC energy.

- 
- [1] V. M. Bannur, J. Phys. G: Nucl. Part. Phys. **32**, 993 (2006).
  - [2] E. V. Shuryak, Phys. Rev. C **80**, 054908 (2009), Erratum-ibid. C **80**, 069902 (2009).
  - [3] STAR Collaboration (John Adams *et al.*), Nucl. Phys. **A757**, 102 (2005); PHENIX Collaboration (K. Adcox *et al.*), Nucl. Phys. **A757**, 184, (2005); B.B. Back *et al.*, Nucl. Phys. **A757**, 28 (2005).
  - [4] H. J. Drescher, A. Dumitru, C. Gombeaud, J. Y. Ollitrault, Phys. Rev. C **76**, 024905 (2007).
  - [5] E. Shuryak, Nucl. Phys. **A774**, 387 (2006).
  - [6] P. Kovtun, D. T. Son, A. O. Starinets, Phys. Rev. Lett. **94**, 111601 (2005).
  - [7] G. Boyd *et al.*, Phys. Rev. Lett. **75**, 4169 (1995); Nucl. Phys. **B 469**, 419 (1996); F. Karsch, Lect. Notes Phys. **583**, 209 (2002); A. Bazavov *et al.*, arXiv:0903.4379.
  - [8] M. Cheng *et al.*, Phys. Rev. D **77**, 014511 (2008) (arXiv:0710.0354).
  - [9] F. Karsch, Lect. Notes Phys. **583**, 209 (2002).
  - [10] R. V. Gavai, Pramana **67**, 885 (2006) (hep-ph/0607050).
  - [11] H. Satz, Nucl. Phys. **A783**, 249 (2007). [arXiv:hep-ph/0609197].
  - [12] E. V. Shuryak, arXiv:hep-ph/0608177.
  - [13] B. A. Gelman, E. V. Shuryak, I. Zahed, Phys. Rev. C **74**, 044908 (2006); Phys. Rev. C **74**, 044909 (2006).
  - [14] E. Shuryak, Nucl. Phys. **A750**, 64 (2005).
  - [15] D. H. Rischke, M. I. Gorenstein, A. Schafer, H. Stocker and W. Greiner Phys. Lett. B **278**, 19 (1992); Z. Phys. **C56**, 325 (1992).
  - [16] B. Sheikholeslami-Sabzevari, Phys. Rev. C **65**, 054904 (2002).
  - [17] P. Arnold and C. Zhai, Phys. Rev. D **50**, 7603 (1994); Phys. Rev. D **51**, 1906 (1995).
  - [18] C. Zhai and B. Kastening, Phys. Rev. **52**, 7232 (1995).
  - [19] K. Kajantie, M. Laine, K. Rummukainen, Y. Schroder, Phys. Rev. D **67**, 105008 (2003).
  - [20] K. Kajantie, M. Laine, K. Rummukainen, Y. Schroder, Phys. Rev. Lett. **86**, 10 (2001).
  - [21] A. Ipp, K. Kajantie, A. Rebhan, A. Vuorinen, Phys. Rev. D **74**, 045016 (2006).
  - [22] P. F. Kelly, Q. Liu, C. Lucchesi and C. Manuel, Phys. Rev. Lett. **72**, 3461 (1994).
  - [23] P. F. Kelly, Q. Liu, C. Lucchesi and C. Manuel, Phys. Rev. D **50**, 4209 (1994).
  - [24] J.-P. Blaizot and E. Iancu, Phys. Rev. Lett. **70**, 3376 (1993).
  - [25] A. Peshier, B. Kampfer, O. P. Pavlenko and G. Soff, Phys. Lett. B **337**, 235 (1994).
  - [26] P. Levai and U. Heinz, Phys. Rev. C **57**, 1879 (1998).
  - [27] A. Peshier, B. Kampfer, O. P. Pavlenko and G. Soff, Phys. Rev. D **54**, 2399 (1996).
  - [28] R. A. Schneider and W. Weise, Phys. Rev. C **64**, 055201 (2001).
  - [29] D. H. Rischke, Prog. Part. Nucl. Phys. **52**, 197 (2004).
  - [30] E. V. Shuryak, Prog. Part. Nucl. Phys. **53**, 273 (2004) [hep-ph/0312227].
  - [31] E. V. Shuryak and I. Zahed, hep-ph/0307267, Phys. Rev. C **70**, 021901 (2004).
  - [32] E. V. Shuryak and I. Zahed, Phys. Rev. D **69**, 014011 (2004) [hep-th/0308073].
  - [33] V. M. Bannur, Phys. Lett. B **362**, 7 (1995).
  - [34] G. Borges, NA50 Collaboration, J. Phys. G **32**, S 381 (2006).
  - [35] A. Adare *et al.* [PHENIX Collaboration], arXiv:0801.0220 [nucl-ex].



- [36] T. Matsui and H. Satz, Phys. Lett. B **178**, 416 (1986).
- [37] See for example: R. Vogt, C. Lourenco and M. J. Leitch, J. Phys. G **34**, S759 (2007); R. Granier de Cassagnac, J. Phys. G **34**, S955 (2007) [arXiv:hep-ph/0701222].
- [38] B. K. Patra and V. J. Menon, Eur. Phys. J. C **37**, 115 (2004), Eur. Phys. J. C **44**, 567 (2005), Eur. Phys. J. C **48**, 207 (2006).
- [39] See for example: R. L. Thews, Nucl. Phys. A **702**, 341 (2002) [arXiv:hep-ph/0111015]; X. Zhao and R. Rapp, Phys. Lett. B **664**, 253 (2008) [arXiv:0712.2407 [hep-ph]].
- [40] A. Pineda and J. Soto, Nucl. Phys. Proc. Suppl. **64**, 428 (1998); N. Brambilla, A. Pineda, J. Soto and A. Vairo, Nucl. Phys. **B566**, 275 (2000).
- [41] W. M. Alberico, A. Beraudo, A. De Pace and A. Molinari, Phys. Rev. D **75**, 074009 (2007).
- [42] D. Cabrera and R. Rapp, Phys. Rev. D **76**, 114506 (2007). [arXiv:hep-ph/0611134].
- [43] Á. Mócsy and P. Petreczky, Phys. Rev. D **77**, 014501 (2008) [arXiv:0705.2559 [hep-ph]].
- [44] N. Brambilla, J. Ghiglieri, A. Vairo and P. Petreczky, arXiv:0804.0993 [hep-ph].
- [45] O. Kaczmarek, F. Karsch, F. Zantow and P. Petreczky, Phys. Rev. D **70**, 074505 (2004) [Erratum-ibid. D **72**, 059903 (2005)]; K. Petrov [RBC-Bielefeld Collaboration], PoS **LAT2006**, 144 (2006) [arXiv:hep-lat/0610041].
- [46] S. Digal, P. Petreczky and H. Satz, Phys. Rev. D **64**, 094015 (2001) [arXiv:hep-ph/0106017].
- [47] C. Y. Wong, Phys. Rev. C **72**, 034906 (2005).
- [48] T. Umeda, K. Nomura and H. Matsufuru, Eur. Phys. J. C **39S1**, 9 (2005); M. Asakawa and T. Hatsuda, Phys. Rev. Lett. **92**, 012001 (2004); G. Aarts et al Phys. Rev. D **76**, 094513 (2007).
- [49] F. Karsch, J Phys: Conference Series **46**, 121 (2006).
- [50] V. Agotiya, V. Chandra and B. K. Patra, Phys. Rev. C **80**, 025210 (2009).
- [51] V. Agotiya, V. Chandra and B. K. Patra, arXiv:nucl-th/0910.0586.
- [52] M. C. Chu and T. Matsui, Phys. Rev. D **37**, 1851 (1988).
- [53] D. Pal, B. K. Patra and D. K. Srivastava, Euro. Phys. J. C **17**, 179 (2000).
- [54] B. K. Patra and D. K. Srivastava, Phys. Lett. B **505**, 113 (2001).
- [55] B. K. Patra, V. Agotiya, and V. Chandra, Eur. Phys. J. C **67**, 465 (2010).
- [56] J. Adams, *et al.*, Phys. Rev. Lett. **92**, 112301 (2004); A. Adare *et al.*, (PHENIX Collaboration), Phys. Rev. Lett. **98**, 232301 (2007).
- [57] F. Karsch, Nucl. Phys. **A698**, 199 (2002); E. Laermann and O. Philipsen, Ann. Rev. Nucl. Part. Sci. **53**, 163 (2003).
- [58] R. A Schneider, Phys. Rev. D **66**, 036003 (2002).
- [59] M. Laine, O. Philipsen, M. Tassler and P. Romatschke, JHEP **03**, 054 (2007).
- [60] A. Beraudo, J. P. Blaizot, C. Ratti, Nucl. Phys. **A806**, 312 (2008).
- [61] Vijai V. Dixit, Modern Physics Letters A **5**, 227 (1990).
- [62] V. M. Bannur, Phys. Lett. **B647**, 271 (2007).
- [63] V. M. Bannur, Eur. Phys. J. C **50**, 629 (2007).
- [64] G. C. Nayak, V. Ravishankar, Phys. Rev. C **58**, 356 (1998); Phys. Rev. D **55**, 6877 (1997).
- [65] R. S. Bhalerao, V. Ravishankar, Phys. Lett. B **409**, 38 (1997).
- [66] A. Jain, V. Ravishankar, Phys. Rev. Lett. **91**, 112301 (2003).
- [67] S. Ichimaru, *Statistical Plasma Physics (Vol. II) - Condensed Plasma* (Addison-Wesley Publishing Company, New York, 1994).
- [68] R. Abe, Progr. Theor. Phys. **21**, 475 (1959).
- [69] S. Ichimaru, Rev. Mod. Phys. **54**, 1017 (1982).
- [70] M. Laine, Y. Schroder, JHEP **0503**, 067 (2005).
- [71] Suzhou Huang, Marcello Lissia, Nucl. Phys. B **438**, 54 (1995).
- [72] E. Braaten and A. Neito, Phys. Rev. D **53**, 3421 (1996) (hep-ph/9510408).
- [73] A. Vuorinen, arXiv:hep-ph/0402242.
- [74] A. Peshier and W. Cassing, Phys. Rev. Lett. **94**, 172301 (2005).
- [75] F. Karsch, E. Laermann and A. Peikert Phys. Lett. **B478**, 447 (2000).
- [76] J. Engels, R. Joswig, F. Karsch, E. Laermann, M. Lutgemeier and B. Petersson, Phys. Lett. B **396**, 210 (1997).
- [77] H. Kouno, M. Maruyama, F. Takagi, and K. Saito, Phys. Rev. D **41**, 2903 (1990).
- [78] A. Muronga, Phys. Rev. Lett. **88**, 062302 (2002).
- [79] F. Karsch, D. Kharzeev and H. Satz, Phys. Lett. B **637**, 75 (2006).
- [80] S. S. Adler et al., (PHENIX Collaboration), Phys. Rev. C **71**, 034908 (2005); S. S. Adler et al., (PHENIX Collaboration), Phys. Rev. C **71** 049901(E) (2005).
- [81] T. Hirano, Phys. Rev. C **65**, 011901(R) (2001); T. Hirano and K. Tsuda, Phys. Rev. C **66**, 054905 (2002).
- [82] K. J. Eskola, K. Kajantie, P. V. Ruuskanen, and K. Tuominen, Nucl. Phys. **B570**, 379 (2000).
- [83] H. Satz, J. Phys. G: Nucl. Part. Phys. **32**, R25 (2006).
- [84] C. Y. Wong, Phys. Rev. C **76**, 014902 (2007).
- [85] S. Datta, F. Karsch, P. Petreczky, and I. Wetzorke, Phys. Rev. D **69**, 094507 (2004).
- [86] Xiao-Ming Xu, D. Kharzeev, H. Satz and Xin-Nian Wang, Phys. Rev. C **53**, 3051 (1996).
- [87] B. K. Patra and V. J. Menon, Eur. Phys. J. C **48**, 207 (2006).
- [88] M. Mishra, C. P. Singh, V. J. Menon and R. K. Dubey, Phys. Lett. B **656**, 45 (2007).
- [89] L. Grandchamp, R. Rapp and G. E. Brown, Phys. Rev. Lett. **92**, 212301 (2004).
- [90] A. Andronic, P. Braun-Munzinger, K. Redlich and J. Stachel, Phys. Lett. B **571**, 36 (2003).
- [91] R. L. Thews and M. L. Mangano, Phys. Rev. C **73**, 014904 (2006); R. L. Thews, Nucl. Phys. A. **783**, 301 (2007).
- [92] S. S. Adler et al., Phys. Rev. Lett. **96**, 032301 (2006).
- [93] A. Adare et al., Phys. Rev. Lett. **97**, 252002 (2006).